On possible applications of the self-channeling in matter of high-power ultrashort laser pulses

A. V. Borovskiĭ, V. V. Korobkin, and A. M. Prokhorov

Institute of General Physics, Russian Academy of Sciences, 38 ul. Vavilova, 117942 Moscow, Russia
(Submitted 5 March 1994)
Zh. Eksp. Teor. Fiz. **106**, 148–160 (July 1994)

In experiments on the self-channeling of a subpicosecond ($\tau=0.6$ ps) high-power ($P\simeq 3\cdot 10^{11}$ W) pulse from a KrF* laser ($\lambda=248$ nm) in gases, intensities in the range $I\simeq 10^{19}$ W/cm^2 have already been obtained in filaments with radius $r\simeq 2\lambda$ and length $L\simeq 2$ mm. Possible applications of the effect are considered: self-compression of pulses, increases in the self-channeled propagation length due to external uptake of energy, new designs for x-ray lasers and harmonic generation, generation of magnetic fields with intensities greater than 10 MG, the possibility of detection of the generation of electron-positron pairs, and the possibility of creating pulsed fast-neutron sources.

1. INTRODUCTION

Recent years, in connection with the vigorous development of pico- and femtosecond laser technology, have seen an acceleration of theoretical and experimental studies of the interaction of ultrashort high-power laser pulses with matter.[1] One of the directions of these investigations has been the nonlinear propagation of ultrashort pulses in matter and new perspectives which have opened up in connection with this.

It is well known that thermal and strictive self-focusing of a light beam, and also the regime of waveguide propagation of light in matter, were predicted by Askar'yan as early as 1962.[2]

At present it is possible to distinguish the following main regimes of nonlinear propagation of laser radiation in matter:

1. The stationary regime

1.1. Waveguide (including pulsed)[2–4] ($L\gg L_d$)

1.2. One or more foci[5,6]

2. The quasi-stationary regime with fast nonlinearity[5,7] $\tau\gg\tau_{nl}$, $\tau\gg L/c$

3. The nonstationary regime with fast nonlinearity[8,9] $\tau\gg\tau_{nl}$, $\tau\ll L/c$

4. The nonstationary regime with slow nonlinearity[10] $\tau\ll\tau_{nl}$.

Here L is the nonlinear propagation length of the laser radiation in the medium, $L_d=d^2/\lambda$ is the diffraction length associated with the transverse dimension d of the pulse in matter, λ is the wavelength of the laser radiation, τ is the duration of the laser pulse, and τ_{nl} is the setup time of the nonlinear response of the medium. Regimes 1, 2, and 4 have already been considered in Refs. 11–13 (reviews) and in Ref. 14.

The present paper considers certain phenomena which arise in the propagation of an ultrashort (usually subpicosecond) high-power laser pulse in a medium with a fast nonlinearity. Under these conditions the model of quasi-stationary self-focusing (case 2) is, generally speaking, inapplicable. For short enough pulses, in describing the dynamics of the pulse, it is impossible to represent the pulse as a set of independent parts, each of which is described by the equation of stationary self-focusing. Each of the various parts of the pulse propagates with its own velocity and interacts with the others in a wavelike fashion. Therefore the given case must be considered a special propagation regime of a high-power ultrashort pulse—the nonstationary regime with fast nonlinearity (case 3).

The above-indicated nonlinear regime of laser pulse propagation was called in Ref. 15 self-channeling of an ultrashort pulse. Its essence consists in the following, namely that a clump of the field plus medium with varying dielectric properties (laser pulse in matter) propagates a distance L which is many times greater than the length of the pulse itself $L_p=c\tau$ and the diffraction length L_d, thus: $L\gg L_p$, L_d. There does not exist a physical channel in the medium along the propagation length of the pulse L. The filament (or channel) here is none other than the trajectory or trace of the motion of a clump of electromagnetic field with small longitudinal and transverse dimensions.

The dynamics of self-channeling of a high-power laser pulse have yet to be studied in detail. A laser pulse, as it propagates, can collapse into a single filament or break up in the transverse direction into several filaments. In turn, the radiation forming the filament can pulsate in a complicated way: it can either stabilize or break up in the longitudinal direction into several clumps with less energy. We emphasize that a rigorous theoretical analysis of these questions must go beyond any approximation tied to the nonlinear Schrödinger equation (NSE),[16,17] and should be based instead on the nonlinear wave equation (NWE).

In Ref. 18 it was predicted that the relativistic-strictive nonlinearity of the electronic component of the plasma leads to self-channeling of an ultrashort high-power laser pulse. The effect was later observed experimentally.[19] The approximate theory of this phenomenon, which makes use of the nonlinear Schrödinger equation, is expounded in Refs. 15, and 20–22, later experimental results are presented in Ref. 23, and the theory which generalizes the nonlinear Schrödinger equation is developed in Ref. 9.

Let us briefly review the physics of the self-channeling effect. A high-power ultrashort laser pulse, as it propagates

 1063-7761/94/070081-07$10.00

in matter, gives rise to a rapid nonlinear ionization of the atoms of the medium. The generation of a plasma decreases the refractive index of the medium, and the forming plasma column possesses linear defocusing properties. However, there exist a number of physical mechanisms leading to a nonlinear increase in the refractive index in the region where the field is intense. These are a) the Kerr effect, which is that high-power laser radiation, by deforming the electron shells of the ions, induces nonlinear dipole moments, b) relativistic increases in the mass of the free electrons oscillating in a strong field; and c) repulsion of the electrons by the ponderomotive force from the region occupied by the field. When the indicated mechanisms exceed a critical power level, they lead to self-channeling of an ultrashort laser pulse. Inside the pulse, a region with reduced electron density forms; that region is shorter than the pulse and moves along with it. The critical power level necessary for relativistic-strictive self-channeling of an ultrashort pulse is[20–22]

$$P_{\rm cr}=1.62\cdot 10^{10}(\omega/\omega_p)^2[W], \tag{1}$$

where $\omega_p^2=4\pi e^2 N_e/m_e$ is the plasma frequency.

Reference 23 reports the observation of self-channeled propagation of a KrF* excimer laser pulse ($\lambda=0.248$ μm, $\tau\simeq 600$ fs, $P\simeq 3\cdot 10^{11}$ W) a distance $L\simeq 2$ mm, which amounts to $\simeq 100L_{\rm d}$. The diameter of the channel was $\simeq 1$ μm. The self-channeled propagation length L was determined by the rate of dissipation of the energy of the pulse in the medium.

It appears that this recently discovered nonlinear propagation regime of ultrashort high-power laser pulses, which leads to strong self-concentration of optical energy in a small region that propagates through the medium, opens up interesting prospects for various new fundamental and practical applications of laser radiation, and also substantially alters the conditions under which some well-known physical processes take place.

In the present paper we discuss not only the peculiarities of some well-known physical effects, but also new effects which can accompany the self-channeled propagation of an ultrashort high-power laser pulse in matter.

2. LONGITUDINAL COMPRESSION OF A LASER PULSE

In the self-channeling of an ultrashort pulse there are two mechanisms altering its temporal shape. First, as the self-channeling regime is entered, the pulse becomes compressed in the transverse coordinate. The effect of this compression on the temporal shape of the pulse in the presence of quasistationary focusing has been discussed in the literature (see, e.g., Refs. 24 and 25). Second, after the self-channeling has been set up, an additional mechanism which alters the longitudinal shape of the pulse comes into play, possibly leading to longitudinal compression. This mechanism is associated with dispersion in the propagation velocities of the various parts of the pulse, which has already been compressed in the transverse direction.

Actually, the group velocity of light propagating in a plasma is close to $v_{\rm g}=c\varepsilon^{1/2}$, where $\varepsilon<1$ is the real part of the dielectric constant. Therefore the central part of the laser pulse propagates with velocity approaching the speed of light, since in the region from which the electrons are repulsed, $\varepsilon\rightarrow 1$. At the same time, the leading edge of the pulse moves somewhat slower. As a result, the pulse becomes compressed. Let us estimate the compression length. The leading edge of the pulse moves a distance L during the time t with the velocity $v_{\rm g}=c(1-\omega_p^2/\omega^2)^{1/2}\simeq c(1-\omega_p^2/2\omega^2)$. At the same time, the central part of the pulse moves a distance $L-L_{\rm p}/2$ with velocity c. Equating times, we obtain the following estimate for the compression length: $L_{\rm c}=L_{\rm p}(\omega^2/\omega_{\rm p}^2-1)$. This estimate is valid, of course, under the condition $\omega^2/\omega_p^2\gg 1$. For example, if $\omega^2/\omega_p^2=100$, then the pulse is compressed at a distance of 100 pulse lengths. For $\tau=500$ fs, $L_{\rm c}=1.5$ cm; for $\tau=100$ fs, $L_{\rm c}=3$ mm. If $\omega^2/\omega_p^2\simeq 10$, then the compression length is decreased by a factor of 10. The above estimates do not take into account the transverse structure of the pulse. Therefore the final conclusion of the possibility of compression and degree of compression can be made only on the basis of a nonstationary and spatially at least two-dimensional theory of self-channeling of an ultrashort high-power laser pulse in matter.

3. INCREASING THE SELF-CHANNELING PULSE PROPAGATION LENGTH VIA AN EXTERNAL SOURCE OF ENERGY

The self-channeling pulse propagation length L is determined by energy losses. The energy loss mechanisms are numerous. Let us list a few of them: ionization and excitation of the atoms and ions of the medium, reverse braking absorption, generation of plasma oscillations, harmonic generation, scattering of radiation by turbulence in the plasma and by the plasma oscillations, partial defocusing of the radiation due to refraction by a nonuniform transverse electron density profile, generation of spontaneous magnetic fields, etc.

An important circumstance in all this is the fact that the walls of the cavitation region that exist inside the traveling pulse can be partially transparent to the laser radiation propagating in it. The reason is that the walls can oscillate with the plasma frequency ω_p, since with the formation of the cavity the electrons are repelled in the transverse direction by the ponderomotive force while the Coulomb attraction force, due to ions remaining in the channel, pulls them back. The plasma oscillations of the repelled electrons (electron cloud) will depend on the time rate of growth of decay of the intensity of the laser radiation in the cavity and on the efficiency of the energy dissipation mechanisms for these oscillations. The laser radiation propagating inside the cavity with its oscillating side walls should partially pass through these walls. Confirmation of the partial transparency of the cavity is provided by experimental observations of scattered laser radiation radiating outward in a narrow cone from the interaction region.[19,23]

The partial transparency of the cavity walls can be used to allow additional energy into the cavity from the region surrounding it. For this purpose, the wavefront of the laser radiation should have a phase component in the form of a cone. Replenishment of the electromagnetic wave inside the cavity might be possible by means of axicon focusing of the laser radiation.[26,27] With external replenishment, one might hope to obtain much longer self-channeling lengths than reported in Ref. 23.

4. X-RAY LASER

The possibility of using laser pulse channeling to amplify x rays was apparently first indicated in Ref. 28. The excitation efficiency of the medium was estimated in Ref. 29. In the self-channeling regime the pulse leaves behind a long narrow filament of plasma consisting of multiply charged ions mainly in excited states and free electrons. In the very near future it may be possible to obtain plasma filaments several centimeters in length, which is sufficient for the generation of stimulated x-ray emission.

At least four different mechanisms of such generation are possible. The first is due to multiphoton excitation of multiply charged ions with subsequent stimulated emission. This mechanism ensures good selectivity of the corresponding levels and has probably been observed in a xenon plasma.[30]

The second mechanism is the so-called self-excitation of the internal electron shells of the ions by collisions with the outer electrons that belong to the same atom and oscillate with relativistic energies sufficient for inelastic processes. This mechanism leads to the appearance of holes in the electron shells. Subsequent dynamics of the holes under a favorable course of events can lead to population inversion of certain transitions. This mechanism was proposed in Ref. 31 and needs further study. The oscillating electrons can lead to the formation of holes not only in the original atom, but also in surrounding atoms since the amplitude of the oscillations can exceed interatomic distances.

The third mechanism is ordinary three-particle recombination of electrons and ions in the plasma filament left behind the pulse. This process usually populates the upper levels of the working ions, and in conjunction with radiative depletion of the lower levels leads to population inversion.[32] Recombinational pumping in the situation under consideration can be efficient as a result of the high density of ions, low plasma temperatures, and optical transparency of the medium in the transverse direction.

A possible fourth mechanism may be optical pumping of the surrounding gas by soft x rays emitted by the plasma filament. In this regard, selective photoionization of subvalence electron shells of heavy atoms in the ambient gas may be promising.[33]

An important feature of the present scheme for exciting the medium during the self-channeling of a laser pulse propagating in a gas is the fact that generation of x rays is possible only inside a band of excited matter remaining behind the pulse. Generation inside the pulse itself is hindered by superstrong dynamic Stark broadening of the excited levels.

A virtue of the scheme is the possibility of amplification of a wave traveling with the speed of light, which is, as is well known, the most efficient regime for amplification of spontaneous emission.

A drawback of the scheme is the excessively small transverse width of the plasma filament, which should lead to large diffraction and refraction losses of the x rays amplified along the filament. All of these losses taken together are referred to as defocusing losses and can be estimated using the following formula:

$$\begin{aligned} &\kappa^- = \sqrt{[(\kappa_1^-)^2 + \sqrt{(\kappa_1^-)^4 + 4(\kappa_2^-)^4}]/2} \\ &\kappa_1^- \simeq 4(\omega_p/\omega d), \\ &\kappa_2^- = [(\kappa^+ - \kappa_a^-)\kappa_d^-]^{1/2}, \\ &\kappa_d^- \simeq \lambda/\pi a^2, \end{aligned} \qquad (2)$$

where κ_1^-, κ^+, κ_a^-, κ_d^- are the coefficients (in units of cm^{-1}) for refraction losses, amplification, absorption, and diffraction losses at the frequency of the amplified radiation $\omega = 2\pi c/\lambda$. Formula (2) was obtained in Ref. 34 on the basis of the theory of the Helmholtz equation for a model Epshteĭn layer, proposed in Ref. 35.

Usually refraction losses predominate. Working with an active medium $L \simeq 1$ cm long, we see that refraction losses place an upper limit on the wavelengths that can be efficiently amplified within the narrow plasma filaments. For efficient amplification, it is desirable to have $\kappa_1^- < L^{-1}$, whereupon $\omega > \omega_p 4L/d$. If, for example, $\omega_p = 10^{15}$ s^{-1}, $L = 1$ cm, $d = 2 \cdot 10^{-3}$ cm, then the wavelength of the amplified radiation should satisfy the inequality $\lambda < 20$ Å.

Let us estimate the energy necessary to obtain a narrow $d = 4\lambda$ string of 10 times ionized plasma. We will assume that an energy $\varepsilon \simeq 2R_8 Z^2$ is expended in the ionization of each atom, where $Z = 10$ and R_∞ is the Rydberg constant. The total energy is then

$$E \simeq \tfrac{1}{4}\pi d^2 L N_i \varepsilon \simeq 8\pi R_\infty \lambda^2 L N_i Z^2. \qquad (3)$$

If $\lambda = 10^{-4}$ cm, $L = 1$ cm, $N_i = 10^{20}$ cm^{-3}, and $Z = 10$, then $E \cong 5 \cdot 10^{-3}$ J/cm.

5. HARMONIC GENERATION

The self-channeled propagation of an ultrashort high-power laser pulse is characterized by certain peculiarities in the generation of harmonics.

There are at least two mechanisms of harmonic generation in a superintense optical field. The first is associated with the excitation of nonlinear dipole moments of the electronic shells of the atoms or ions[36,37] and leads to the appearance of a large number of odd harmonics of comparable amplitude. The second mechanism is due to the emission of free electrons moving in the superintense field along complicated trajectories in the form of figure eights or spirals.[38,39]

For efficient generation of harmonics it is necessary that the phase velocities of electromagnetic wave propagation in a prescribed direction for the fundamental frequency and the frequency of the harmonic coincide (the condition of phase synchronism). If this condition is not fulfilled, then the intensity of the harmonic oscillates in space at a low level in comparison with that of the fundamental. In the direction of propagation of the fundamental the spatial period of the oscillations of the harmonic is $L_s = \lambda n / 2\Delta n$, where Δn is the difference in the refractive index at the frequencies considered.

A peculiarity of self-channeling is that the diffraction length for the harmonic $L_{Dg} = d^2/\lambda_g$ for $d \simeq 4\lambda$ turns out to be less than the phase synchronism length L_s. This means that the harmonic radiation is drawn out of the channel through its lateral surface into the diffraction cone. In the far zone the distribution of harmonic radiation is determined by interference of the diffracted fields that leak out at various points along the channel.

Indeed, let us consider the fundamental radiation propagating along the z axis in the self-channeling regime between the points $z=0$ and $z=L$. We assume that the radiation propagates without attenuation, and the field has a Gaussian distribution and a planar wavefront in any transverse cross section of the channel. Then the expression for $E(r,z,t)$ of the fundamental radiation has the form

$$E(r,z,t) = E_0 \exp(-r^2/d_0^2) \exp[i(\omega t - kz)], \tag{4}$$

where d_0 is the transverse radius of the beam.

Let us consider some fixed $z=\xi$. We assume that at each point of the transverse cross section there takes place an in-phase conversion of the fundamental radiation into a harmonic of multiplicity m with amplitude coefficient γ_m. The transverse distribution of the real amplitude of the field of the harmonic generated at $z=\xi$ in the interval $d\xi$ is

$$d|E_g| = \gamma_m [E_0 \exp(-r^2/d_0^2)]^m d\xi \equiv \exp(-r^2/d_{0g}^2) dE_{0g}, \tag{5}$$

and hence the increment of the real amplitude of the field of the harmonic along the axis is $dE_{0g} = \gamma_m E_0^m d\xi$, and the transverse radius of the harmonic is $d_{0g} = d_0/m^{1/2}$.

The element of the field of the harmonic generated at $z=\xi$ has a planar wavefront and propagates in space according to the laws for a Gaussian beam

$$dE_g(r,z-\xi,t) = dE_{0g} \frac{d_{0g}}{d_g(z-\xi)} \exp\left[-\frac{r^2}{d_g^2(z-\xi)}\right] \times \exp\left\{i\left[\omega_g t - k_g z + k_g \xi + \varphi_g(\xi) + \alpha_g(z-\xi) - \frac{k_g r^2}{2R_g(z-\xi)}\right]\right\}, \tag{6}$$

where d_{0g} is the minimum beam radius (the radius at the beam waist), $d_g^2(z) = d_{0g}^2(1 + z^2/b_g^2)$, $\alpha_g(z) = \arctan(z/b_g)$, $R_g(z) = z(1 + z^2/b_g^2)$, $b = \pi d_{0g}^2 n_g/\lambda_g$, $k_g = 2\pi n_g/\lambda_g$, and n_g is the refractive index. In formula (6) the term $\varphi_g(\xi)$ is the contribution to the phase of the harmonic due to propagation of the pump wave along the z axis and is equal to $\varphi_g(\xi) = -k\xi\omega_g/\omega$.

The total magnitude of the harmonic field is obtained by integrating expression (6) with respect to ξ over the interval $0 < \xi < L_p$. We denote the look angle by θ. Then $\tan\theta = r/z$, and we obtain

$$E_g(z,\theta,t) = \gamma_m E_0^m \frac{b_g}{z} \exp\left(-\frac{b_g^2}{d_{0g}^2} \tan^2\theta\right) \times \exp\left\{i\left[\omega_g t - k_g z\left(1 + \frac{1}{2}\tan^2\theta\right) + \pi/2\right]\right\} \times [\exp(iL_p\Delta k) - 1]/i\Delta k L_p, \tag{7}$$

where the dephasing Δk is equal to

$$\Delta k = k_g(1 - \tfrac{1}{2}\tan^2\theta) - \frac{\omega_g}{\omega} k. \tag{8}$$

Thus, efficient harmonic generation should be observed at angles for which the condition $\tan\theta = [2(1 - n/n_g)]^{1/2}$ is fulfilled. Here $n < n_g$. In addition, the angle θ should lie inside the diffraction cone of the harmonic: $\theta < \theta_d = \pi d_{0g}^2/\lambda_g$.

Earlier we indicated two mechanisms of harmonic generation in strong fields. These mechanisms differ in the shape of the transverse spatial zones of generation. In the case of harmonic generation by the dipole moments of the electronic shells of the atoms and ions of the medium the transverse generation zone is a disk since the process takes place inside the cavitation channel, inside which the radiation intensity is maximum and there are no free electrons. In the case of harmonic generation by nonlinear currents of free electrons, the generation zone is a ring, since the free electrons are expelled from the channel and the process takes place inside the electron cloud surrounding the cavitation channel. The various transverse zones have different angular intensity distributions of the interference rings of the harmonics in the far zone.

Recording of the harmonics can serve as a method of determining the phase velocity of the fundamental wave, the phase synchronism lengths, and the power captured into the self-channeling regime, and also enables one to distinguish the different harmonic generation mechanisms.

6. GENERATION OF STRONG MAGNETIC FIELDS

The self-channeling regime is accompanied by the action of powerful electromagnetic forces on each electron located in the interaction zone. As a rule, the motion of the electron can be broken down into two components: high-frequency oscillations and a drift motion.

Let us first discuss the drift motion of the electrons. At the leading edge of the laser pulse the electron absorbs a large number of photons from the electromagnetic field and receives momentum in the direction of propagation of the field. The ponderomotive force repels the electron in the transverse direction.

The equation for the transverse drift of the electron describes the translation of an electron made heavier by the relativistic effect in the presence of high-frequency oscillations, in some force field which communicates to the electron an amount of kinetic energy equal to the difference in the ponderomotive potential between the initial and final points of its translation. The magnitude of this potential in the relativistic case is[40]

$$\varphi_{\rm po}=m_{\rm e,0}c^2\gamma,\quad \gamma=(1+\alpha)^{1/2}, \tag{9}$$

where $\alpha=I/I_r$. Here I is the intensity and I_r is a parameter which is called the relativistic intensity

$$I_{\rm r}=\frac{m_{\rm e,0}^2\omega^2c^3}{4\pi e^2}=2.75\cdot 10^{18}\left(\frac{1}{\lambda[\mu{\rm m}]}\right)^2 {\rm W/cm}^2. \tag{10}$$

Formula (10) was derived for the case of circularly polarized radiation.

The longitudinal motion of the electron differs from its transverse motion. The corresponding equation is discussed, e.g., in Ref. 41.

It is easy to show by a simple estimate that an electron repelled by the ponderomotive force can acquire a kinetic energy in excess of 1 MeV. The velocity of such an electron approaches the speed of light. Once the electron leaves the strong field region, a Coulomb force from the positive ion swarm begins to act on it pulling it back. The ions remain immobile for some time as a result of their large inertia. The Coulomb force retards the repelled electron. The competition between the repulsive ponderomotive force and the attractive Coulomb force gives rise then to oscillations at the plasma frequency of the electrons, which form a cloud around the strong field zone. At the trailing edge of the pulse, the electrons return into the ion swarm, neutralizing it, and thermalization of the plasma oscillations takes place.

Let us estimate the magnitude of the magnetic field which arises as a result of the repulsion of the electrons from the strong high-frequency field region. The circulation of the magnetic field about a closed contour is equal to the total current through the surface bounded by this contour. Let the contour be a circle in the plane perpendicular to the direction of propagation of the laser pulse and centered on the axis of the beam, which we take to be axially symmetric. We choose the radius of the circle to be equal to the transverse dimension of the laser pulse $R=L_\perp$. We then have

$$B\cdot 2\pi R=\frac{4\pi}{c}\int_S jdS\simeq\frac{4\pi}{c}en_ec\pi R^2.$$

We obtain an overestimate of the magnitude of the magnetic field

$$B\simeq 2\pi en_eL_\perp . \tag{11}$$

Assuming, for example, that a pulse from a neodymium laser propagates in the medium in the self-channeled regime, and that $L_\perp\simeq 2\lambda=2\cdot 10^{-4}$ cm and $n_e\simeq 10^{20}$ cm^{-3}, we obtain $B\simeq 60$ MG.

Magnetic fields of the indicated magnitude should cardinally influence the properties of the medium and the picture of nonlinear propagation of an ultrashort laser pulse in the medium.

Numerical modeling of the action of a 100-fs pulse with an intensity of 10^{19} W/cm^2 on a solid-state target, reported in Ref. 42, gave a magnetic field intensity B in the neighborhood of 250 MG.

Let us consider in more detail the high-frequency component of the motion of the electrons. The case of a circularly polarized electromagnetic field is of interest. The electrons, turning in micro-circles in such a transversely inhomogeneous field, create an uncompensated circular current around the axis of propagation of the laser beam. Let us estimate the magnitude of the magnetic field created by this current parallel to the axis. We take as the integration contour a narrow band of width Δz lying in a plane passing through the beam axis. The lateral legs of the contour are perpendicular to the beam axis and extend to infinity. The circulation of the magnetic field about the above-indicated contour is equal to $B\cdot\Delta z$, since the field lines are perpendicular to the lateral legs of the contour and the field vanishes at infinity. The contribution to the quasistatic current flowing through the band defined by the contour comes from electrons circling about the length element of the contour Δz. The electrons, turning in circles which cut through the surface of the band twice give zero contribution to the total current. As a result, we obtain the estimate

$$B\Delta z=\frac{4\pi}{c}\frac{ev}{2\pi r}n_e\pi r^2\Delta z.$$

Next, applying the relativistic formulas for the velocity of the electron v and the radius of the circle r about which it turns in the circular field with intensity I, we obtain the following formula for the magnetic field:

$$B=2\pi en_ek^{-1}\cdot\frac{I}{I+I_r}. \tag{12}$$

Here $k=\omega/c$ is the wave number, and I_r is the relativistic intensity. In the ultrarelativistic limit $I\gg I_r$ for a neodymium laser and $n_e=5\cdot 10^{20}$ cm^{-3}, the magnitude of the field $B\simeq 24$ MG.

Note the following fact. A plasma in which the electrons revolve synchronously in micro-circles is a ferromagnet. Estimating the magnetic field in terms of the magnetization of the medium $\mathbf{B}=4\pi\mathbf{P}$, where $\mathbf{P}=n_e\mathbf{p}$, and $\mathbf{p}$ is the magnetic moment of an individual revolving electron, leads to the same formula (12).

7. GENERATION OF ELECTRON-POSITRON PAIRS

The practical realization of relativistic-strictive self-channeling of an ultrashort high-power laser pulse in gases[19,23] opens up possibilities for answering the question of the generation of electron-positron pairs in the collisions of electrons oscillating in the field of an intense optical wave with the nuclei of heavy elements.

Quantum electrodynamics, as is well known, predicts the possibility of creating electron-positron pairs in the field of an intense electromagnetic wave. The probability of pair creation at the focus of an ideal lens for focusing coherent laser radiation was first estimated in Refs. 43 and 44. The most favorable results were obtained for pair generation by electrons scattering off the nuclei of heavy elements in the presence of an intense electromagnetic wave.[46,47]

Pair generation will take place more efficiently in a circularly polarized field. An electron in such a field, whose intensity varies slowly in time, moves in a spiral trajectory and much of the time possesses the energy needed for pair generation, in comparison with the case of a linearly polarized field.

Let us estimate the energy of a revolving electron:

$$E=(p^2c^2+m_{e,0}^2c^4)^{1/2}=m_{e,0}c^2(1+\alpha)^{1/2},\quad \alpha=I/I_r. \tag{13}$$

The energy E should exceed the sum, first of all, of the rest masses of two electrons and a positron $3m_{e,0}c^2$, which obtains for $\alpha>8$; second, roughly 10 bond energies of the electrons with the heavy atom. The latter is necessary to ensure that the relativistic electron will break through the electron shells of the atom. Taking all this into account, the estimate for the threshold energy for $Z\simeq 90$ is

$$I>I'=1.1\cdot 10^{20}\left(\frac{1}{\lambda[\mu\mathrm{m}]}\right)^2 \mathrm{W/cm^2}. \tag{14}$$

In the setup of the experiment to observe electron-positron pair generation, self-channeled propagation of an ultrashort laser pulse in a gas is of no less interest than focusing on a solid-state target.

Indeed, the pair yield in both cases can be estimated in the following way:

$$N\simeq n_e n_a\langle\sigma v_e\rangle V\tau\simeq n_e n_a \sigma c V\tau. \tag{15}$$

Here n_a is the concentration of heavy target atoms, σ is the cross section of the process, V is the volume of the interaction region of the radiation with the medium, and τ is the duration of the laser pulse. The relative yield of reaction products for the cases of self-channeling in a gas and focusing on a solid-state target is

$$N_g/N_t=(\rho_g^2 V_g\sigma_g)/(\rho_t^2 V_t\sigma_t) \tag{16}$$

i.e., the ratio of the squares of the densities of the two media, the volumes of the interaction regions, and the cross sections. For a neodymium laser, the first ratio ρ_g^2/ρ_t^2 can be equal to $\simeq 10^{-4}$. The interaction region in the case of self-channeling, according to Refs. 20–22, is a tubular shell with thickness $\Delta r\simeq\lambda$ and outer radius $\simeq 3\lambda$. Therefore $V_g\simeq 5\pi\lambda^2 L$, where L is the pulse propagation length in the medium. The interaction region in the solid-state target is a disk of thickness $\simeq\lambda$ and radius $\simeq 3\lambda$. The radius is determined by the quality of focusing of the system. Therefore $V_t\simeq 9\pi\lambda^2$. The ratio of volumes V_g/V_t can be equal to $\simeq 0.5L/\lambda\simeq 0.5\cdot 10^4$ for lengths $L\simeq 1$ cm. In the near-threshold region, the cross section of the process grows abruptly with growth of the energy of the oscillating electron, i.e., with increase of the intensity of the radiation in the interaction region. It would seem that higher radiation intensities can be realized in the case of self-channeling of an ultrashort pulse in a gas than by focusing on a solid-state target. Thus, self-channeling of laser radiation in a gas is promising for electron-positron pair generation.

In conclusion, let us estimate the number of head-on collisions of the relativistic electrons with the target nuclei in the self-channeling experiment. Taking the cross section to be equal to the transverse area of the heavy nucleus $\sigma=\pi r_0^2 A^{2/3}=5.5\cdot 10^{-25}$ cm^2, where $r_0=1.2\cdot 10^{-13}$ cm, $A=238$; e.g., for $n_e=10^{21}$ cm^{-3}, $n_a=10^{19}$ cm^{-3}, $\lambda=1.06\cdot 10^{-4}$ cm, $L=1$ cm, $\tau=10^{-12}$ s, we obtain $N_g=3\cdot 10^7$. The pair yield will, of course, be less than this.

8. PULSED NEUTRON SOURCE

It is of interest to realize the regime of self-channeling of an ultrashort laser pulse in a gas mixture of deuterium and tritium. The ponderomotive forces expelling the electrons from the region occupied by the strong field lead to the formation of a positively charged swarm of ions with transverse dimension $d\simeq 4\lambda$. The Coulomb forces then push the positively charged ion swarm apart, transferring to the ions a large kinetic energy. The ion swarm simply explodes. After the impact of the scattering ions with the adjacent layers of gas, the ions thermalize and thermonuclear reactions can take place in this region at times of the subsequent gas-dynamic scattering of the plasma. Let us make some estimates.

For simplicity, let us consider the Coulomb-repulsive disintegration of an extended ionic cylinder of length l and radius $R=2\lambda$ $(l\gg R)$. The mean Coulomb energy per ion in such a cylinder is

$$\overline{W^{(\mathrm{cyl})}}=\frac{\pi}{2}e^2NR^2. \tag{17}$$

The characteristic time of conversion of Coulomb potential energy into kinetic energy of directed motion of the ions $\tau_q\simeq\omega_{p,i}^{-1}$, where $\omega_{p,i}^2=4\pi e^2N/m_i$ is the ionic plasma frequency. If $\omega_{p,i}=10^{13}$ s^{-1}, then $\tau_q=100$ fs. This means that if we use laser pulses with duration $1>\tau>0.1$ ps, then after such a pulse has acted, all of the Coulomb energy of the ions in the channel will have been converted into the kinetic energy of the explosive motion of the ions. A large part of this energy is converted into heat during the thermalization stage.

We estimate the neutron yield from the formula

$$N\simeq\tfrac{1}{4}N^2\langle\sigma v\rangle t_0 V. \tag{18}$$

Here the gas-dynamic scattering time of the heated ion swarm $t_0\simeq R/v$, where v is the thermal velocity. We obtain $N\simeq(\pi/4)N_i^2\sigma R^3 L$. Let us make some estimates for the following parameter values: $l=c\tau\simeq 3\cdot 10^{-2}$ cm, $R=2\lambda\simeq 2\cdot 10^{-4}$ cm, $N_i\simeq 10^{21}$ cm^{-3}, $L\simeq 10^{-1}$ cm. For the mean Coulomb energy we obtain $\overline{W^{(\mathrm{cyl})}}\simeq 9\cdot 10^6$ eV, $\sigma\simeq 10^{-25}$ cm^2. For the neutron yield, $N\simeq 6\cdot 10^4$. For the power of the neutron pulse, $P=Nc/L\simeq 6\cdot 10^{16}$ s^{-1}. The neutron yield can be increased using a shorter-wavelength laser, since $N\simeq N_{\mathrm{sh}}\propto\omega^2$, but $R\propto\omega^{-1}$, so $N\propto\omega$. The neutron yield is probably somewhat higher than the estimate given here due to additional compression of the ions due to impact with the adjacent layer of gas.

Thus, self-channeled propagation of an ultrashort high-power laser pulse in matter opens up a wide range of possibilities for the investigation of new physical phenomena in superstrong optical fields, and also substantially changes the form of well-known phenomena. An important attendant circumstance here is the conversion of the energy of the electromagnetic field into the energy of plasma oscillations. Effects described in this paper are preserved with some changes of form in this case. It would seem promising to us to carry out additional experimental and theoretical studies of the directions enumerated in this paper.

The authors thank G. A. Askar'yan and Ch. K. Mukhtarov for valuable remarks.

This work was supported by the Russian Fund for Fundamental Research, Grant No. 93-02-14328.

[1] B. Luther-Davies, E. G. Gamaliĭ, Y. Wang, A. V. Rode, and V. T. Tikhonchuk, Kvantovaya Electron. **19**, 317 (1992) [Sov. Phys. Quantum Electron. **22**, 289 (1992)].

[2] G. A. Askar'yan, Zh. Eksp. Teor. Fiz. **42**, 1567 (1962) [Sov. Phys. JETP **15**, 1088 (1962)].

[3] V. I. Talanov, Izv. Vyssh. Uchebn. Zaved., Radiofiz. **7**, 564 (1964).

[4] V. E. Zakharov, V. V. Sobol'ev, and V. S. Synakh, Zh. Eksp. Teor. Fiz. **60**, 136 (1971) [Sov. Phys. JETP **33**, 77 (1971)].

[5] A. L. Dyshko, V. N. Lugovoĭ, and A. M. Prokhorov, Pis'ma Zh. Eksp. Teor. Fiz. **6**, 655 (1967) [JETP Lett. **6**, 146 (1967)].

[6] A. A. Amosov, N. S. Bakhvalov, Ya. M. Zhileĭkin *et al.*, Pis'ma Zh. Eksp. Teor. Fiz. **30**, 119 (1979) [JETP Lett. **30**, 108 (1979)].

[7] V. V. Korobkin, A. M. Prokhorov, R. V. Serov, and M. Ya. Shchelov, Pis'ma Zh. Eksp. Teor. Fiz. **11**, 153 (1970) [JETP Lett. **11**, 94 (1970)].

[8] A. G. Litvak and V. I. Talanov, Izv. Vyssh. Uchebn. Zaved., Radiofiz. **10**, 539 (1967).

[9] A. V. Borovskiĭ and A. L. Galkin, Zh. Eksp. Teor. Fiz. **104**, 3311 (1993) [Sov. Phys. JETP **77**, 562 (1993)].

[10] S. A. Akhmanov, S. P. Sukhorukov, and R. V. Khokhlov, Zh. Eksp. Teor. Fiz. **51**, 290 (1966) [Sov. Phys. JETP **24**, 431 (1966)].

[11] G. A. Askar'yan, Usp. Fiz. Nauk **107**, 507 (1972) [Sov. Phys. Usp. **15**, 517 (1972–1973)]; Usp. Fiz. Nauk **111**, 249 (1973) [Sov. Phys. Usp. **16**, 680 (1973)].

[12] V. V. Korobkin, Usp. Fiz. Nauk **107**, 512 (1972) [*sic*].

[13] V. I. Lugovoĭ and A. M. Prokhorov, Usp. Fiz. Nauk **111**, 203 (1973) [Sov. Phys. Usp. **16**, 658 (1973)].

[14] S. A. Akhmanov, V. A. Visloukh, and A. S. Chirkin, *Optics of Femtosecond Laser Pulses* [in Russian], Nauka, Moscow (1988).

[15] A. B. Borisov, A. V. Borovskiĭ, V. V. Korobkin, A. M. Prokhorov, C. K. Rhodes, O. B. Shiryaev, Zh. Eksp. Teor. Fiz. **101**, 1132 (1992) [Sov. Phys. JETP **74**, 604 (1992)].

[16] V. I. Talanov, Pis'ma Zh. Eksp. Teor. Fiz. **2**, 218 (1965) [JETP Lett. **2**, 138 (1965)].

[17] G. Z. Sun, E. Ott, Y. C. Lee, and P. Guzdar, Phys. Fluids **30**, 526 (1987).

[18] A. B. Borisov, A. V. Borovskiĭ, V. V. Korobkin, A. M. Prokhorov, C. K. Rhodes, and O. B. Shiryaev, Kratkie Soobshcheniya po Fizike, No. 9, pp. 3–7 (1991).

[19] A. B. Borisov, A. V. Borovskii (Borovskiy), V. V. Korobkin, A. M. Prokhorov, O. B. Shiryaev, X. M. Shi, T. S. Luk, A. McPherson, J. L. Solem, K. Boer, and C. K. Rhodes, Phys. Rev. Lett. **68**, 2309 (1992).

[20] A. B. Borisov, A. V. Borovskii (Borovskiy), O. B. Shiryaev, V. V. Korobkin, A. M. Prokhorov, J. L. Solem, T. S. Luk, K. Boer, and C. K. Rhodes, Phys. Rev. A **45**, 5830 (1992).

[21] A. B. Borisov, A. V. Borovskii (Borovskiy), O. B. Shiryaev, V. V. Korobkin, A. M. Prokhorov, and C. K. Rhodes, Laser Physics **1**, 546 (1991).

[22] A. B. Borisov, A. V. Borovskiĭ, C. K. Rhodes, and O. B. Shiryaev, Trudy IOFAN **41**, 3 (1993).

[23] A. B. Borisov, A. V. Borovskii (Borovskiy), V. B. Karpov, V. V. Korobkin, A. M. Prokhorov, O. B. Shiryaev, X. M. Shi, T. S. Luk, A. McPherson, J. L. Boer, and C. K. Rhodes, *Proceedings of the Third International Symposium on X-Ray Lasers*, Munich, May 1992.

[24] R. G. Brewer, J. R. Lifsitz, E. Garmire *et al.*, Phys. Rev. **166**, 326 (1968).

[25] T. K. Gustafson, J. P. Taran, H. A. Haus *et al.*, Phys. Rev. **177**, 306 (1969).

[26] F. V. Bunkin, V. V. Korobkin, Yu. A. Kurinyĭ, L. Ya. Polonskiĭ, and L. N. Pyatnitskiĭ, Kvantovaya Elektron. **10**, 443 (1983) [Sov. Phys. Quantum Electron. **13**, 254 (1983)].

[27] V. V. Korobkin, M. Yu. Marin, V. I. Pil'skiĭ, L. Ya. Polonskiĭ, and L. N. Pyatnitskiĭ, Kvantovaya Elektron. **12**, 959 (1985) [Sov. Phys. Quantum Electron. **15**, 631 (1985)].

[28] J. C. Solem, T. S. Luk, K. Boyer, and C. K. Rhodes, IEEE J. Quantum Electron. **25**, 2423 (1989).

[29] A. B. Borisov, A. V. Borovskii (Borovskiy), V. V. Korobkin, O. B. Shiryaev, X. M. Shi, T. S. Luk, A. McPherson, J. L. Solem, K. Boer, and C. K. Rhodes, *Proceedings of the International Conference on Lasers-91*, San Francisco, December 9–13, 1991.

[30] A. McPherson, G. Gibson, H. Jara, U. Johann, T. S. Luk, I. A. McIntyre, K. Boer, C. K. Rhodes, J. Opt. Soc. Am. B **4**, 595 (1987).

[31] V. V. Korobkin and M. Yu. Romanovskiĭ, Pis'ma Zh. Eksp. Teor. Fiz. **53**, 493 (1991) [JETP Lett. **53**, 470 (1991)].

[32] A. V. Borovskii (Borovskiy), E. V. Chizhonkov, A. L. Galkin, and V. V. Korobkin, Appl. Phys. B **50**, 297 (1990).

[33] A. V. Borovskiĭ, P. A. Golovinskiĭ, S. I. Derzhavin, B. V. Mokrov, A. A. Sirotkin, and V. S. Rostovtsev, "Photoionization pumping of laser media," in *Nonlinear Spectroscopy of Atomic and Diatomic Molecules*, N. B. Delone (ed.), Moscow (1988), p. 200.

[34] A. V. Borovskiĭ, V. V. Korobkin, and Ch. K. Mukhtarov, Kvantovaya Elektron. **14**, 2269 (1987) [Sov. Phys. Quantum Electron. **17**, 1447 (1987)].

[35] F. V. Bunkin and V. P. Bykov, Kvantovaya Elektron. **13**, 869 (1986) [Sov. Phys. Quantum Electron. **16**, 1323 (1986)].

[36] M. Ferray, A. L'Huillier, X. F. Li, L. A. Lompre, G. Mainfray, and C. Manus, J. Phys. B **21**, L31 (1988).

[37] J. H. Eberly, Q. Su, and J. Javanainen, Phys. Rev. Lett. **62**, 881 (1989).

[38] P. Sprangle, E. Esarey, and A. Ting, Phys. Rev. Lett. **64**, 2011 (1990); Phys. Rev. A **41**, 4463 (1990); Phys. Fluids B **2**, 1390 (1990).

[39] J. M. Rax and N. J. Fisch, Phys. Rev. Lett. **69**, 772 (1992).

[40] A. I. Akhiezer and R. V. Polovin, Zh. Eksp. Teor. Fiz. **30**, 915 (1956).

[41] S. P. Goreslavskii, N. B. Narozhnyi, and V. P. Yakovlev, Laser Physics **1**, 670 (1991).

[42] S. C. Wilks, W. L. Kruer, M. Tabak, and A. B. Langdon, Phys. Rev. Lett. **69**, 1383 (1992).

[43] F. V. Bunkin and A. M. Prokhorov, *The Interaction of Electrons with High-Intensity Optical Radiation, Commemorative Volume Honoring Prof. Kastler*, Paris (1968).

[44] F. V. Bunkin and I. I. Tugov, Dokl. Akad. Nauk **187**, 541 (1969) [Sov. Phys. Dokl. **187**, 678 (1970)].

[45] V. I. Ritus, Doctoral Dissertation, Lebedev Physical Institute, Moscow (1969).

[46] F. V. Bunkin and A. M. Prokhorov, *Polarization de la Matière et du Rayonnement, Volume Jubilaire en L'Honneur d'Alfred Kastler*, Paris (1969).

[47] F. V. Bunkin and A. E. Kazakov, Dokl. Akad. Nauk **193**, 1274 (1970) [Sov. Phys. Dokl. **193**, 758 (1971)].

Translated by Paul F. Schippnick